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# Some Conjectures on the Limit of Infinite Higgs Mass

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**Abstract**

We consider a possible field theory candidate for the electroweak $SU(2)\otimes U(1)$ model where the limit of infinitely sharp Higgs potential is performed. We show that it is possible to formulate such a limit as a Stückelberg massive non Abelian gauge theory.

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# 1 Introduction

The electroweak model [1] as a part of the Standard Model has had impressive confirmations [2] and great is the expectation for the experimental evidence of the last important part: the Higgs sector. The confirmation of the spontaneous breaking of the $SU(2)\otimes U(1)$ via Higgs mechanism would be an important step in the development of quantum field theory. Gauge theories have proved to be the correct framework for the description of the elementary particle world. Their implementation by means of the spontaneously broken symmetry mechanism is then expected as a further development.

This satisfactory situation is not a reason to stop searching for new implementations of gauge theories. In particular we are interested in the formulation of massive gauge theories à la Stückelberg, where the mass for the gauge field is introduced without breaking the local gauge invariance (BRST invariance in the quantum version of the theory). Of course we face a presently insurmountable objection on the non-renormalizability of the perturbation theory. However many options could in principle be at disposal as a way out to this objection, as lattice calculation of non perturbative effects.

In the present paper we look for a formulation of the electroweak $SU(2)\otimes U(1)$ model where the masses of the vector mesons are introduced via the Stückelberg method. The existence of such a formulation is important for two reasons. First it opens the way to a possible use of the equivalence theorem in order to set a bridge with the unitary gauge of massive gauge theory. Second it would allow topological arguments on the ew model based on the fact that the Stückelberg mass makes use of the flat connection built on a non linear sigma model.

We imagine our final model as a limit for infinitely sharp Higgs potential (in standard notations $\lambda \to \infty$) without, however, committing ourselves with a limit on the mass of the Higgs boson. The two problems seem to us very distinct and only a reasonable theory of the non-linear sigma model would be able to shed some light on their inter correlations.

The formal limit of $\lambda \to \infty$ has been considered by some authors [3] and there is unanimous consensus that the boson sector is described by a non-linear sigma model. In the present paper we show that the scalar field can be accommodated so that it appears only via a flat connection, i.e. a pure-gauge field.

The plan of the paper is somehow reversed. We prefer to construct a model which has no physical interpretation, but serves to us to put down the rational for the construction of the model we are really interested in. Our guide in the construction will simply be the compatibility of the main features of the present phenomenology with the tree approximation of the theory.

## 2 Preliminary considerations

Let us consider a fully globally symmetric $SU(2)$ with a mass term in the Proca gauge

$$S = \int d^4x \left( -\frac{1}{4} G_{a\mu\nu} G_a^{\mu\nu} + m^2\ Tr\ [A_\mu A^\mu] \right), \tag{1}$$

where

$$A_\mu = \frac{1}{2}\tau_a A_{a\mu}. \tag{2}$$

$\tau_a$ are the Pauli matrices. Let us now perform a formal operator valued local $SU(2)$ transformation in order to introduce the Stückelberg field

$$A'_\mu = \Omega^\dagger A_\mu \Omega - \frac{i}{g}\Omega^\dagger \partial_\mu \Omega, \tag{3}$$

with the constraint

$$\Omega \in SU(2) \quad \Longrightarrow \quad \Omega^\dagger = \Omega^{-1}, \quad \det\Omega = 1. \tag{4}$$

One gets

$$\begin{aligned} S = \int d^4x \Big\{ & -\frac{1}{4} G_{a\mu\nu} G_a^{\mu\nu} \\ & + \frac{m^2}{g^2}\ Tr\ \Big[ \Big( g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega \Big) \Big( g\Omega^\dagger A^\mu \Omega - i\Omega^\dagger \partial^\mu \Omega \Big) \Big] \Big\}. \end{aligned} \tag{5}$$

**Proposition 1:** *Each element of the matrix (each term in round brackets in eq. (5))*

$$\Big( g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega \Big)_{ab} \tag{6}$$

*is invariant under local $SU(2)$ left transformations*

$$\begin{aligned} A'_\mu &= U_L A_\mu U_L^\dagger - \frac{i}{g} U_L \partial_\mu U_L^\dagger \\ \Omega' &= U_L \Omega. \end{aligned} \tag{7}$$

In fact

$$\begin{aligned} & g\Omega'^\dagger A'_\mu \Omega' - i\Omega'^\dagger \partial_\mu \Omega' \\ & = g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger U_L^\dagger U_L \partial_\mu U_L^\dagger U_L \Omega - i\Omega^\dagger U_L^\dagger \partial^\mu (U_L \Omega) \\ & = g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial^\mu \Omega. \end{aligned} \tag{8}$$

As a consequence of the above result, we can construct an arbitrary number of invariants since there are no constraints on the dimensionality of the coupling constants. Among them some are bilinear in the gauge field. For instance the following two terms

$$\int d^4x \frac{m'^2}{g^2} v^\dagger \Big[ \Big( g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega \Big)^2 \Big] v, \tag{9}$$

$$\int d^4x \frac{m'^2}{g^2} \Big( v^\dagger \Big( g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega \Big) v \Big)^2, \tag{10}$$

where $v$ is any constant spinor which we choose to normalize by

$$v^\dagger v = 1. \tag{11}$$

The proposition 1. can be extended to $SU(2) \otimes U(1)$. One enlarges the local gauge transformations in eq. (7) by introducing a further abelian gauge field [5]

$$\begin{aligned} B'_\mu &= B_\mu - \frac{1}{g'} \partial_\mu \lambda \\ A'_\mu &= U_L A_\mu U_L^\dagger - \frac{i}{g} U_L \partial_\mu U_L^\dagger \\ \widehat{\Omega}' &= e^{i\lambda} U_L \widehat{\Omega} \\ U_L &\in SU(2). \end{aligned} \tag{12}$$

The transformation properties of $\widehat{\Omega}$ implies that four real fields are necessary in order to describe the degree of freedom

$$\widehat{\Omega} = \exp(i\phi_B)(\phi_0 + i\tau_a \phi_a), \qquad \phi_0^2 + \vec{\phi}^2 = 1. \tag{13}$$

**Proposition 2:** *Each element of the following matrix*

$$\Big( g' B_\mu + g\widehat{\Omega}^\dagger A_\mu \widehat{\Omega} - i\widehat{\Omega}^\dagger \partial_\mu \widehat{\Omega} \Big)_{ab} \tag{14}$$

*is invariant under the transformations (12).* Thus, again, any Lorentz invariant function of (14) is a possible term of the action. In particular the action for eq. (9) becomes

$$\begin{aligned} S = \int d^4x \Big( &- \frac{1}{4} G_{a\mu\nu} G_a^{\mu\nu} - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} \\ &+ \frac{m^2}{g^2}\, Tr\, \Big\{ \Big( g' B_\mu + g\widehat{\Omega}^\dagger A_\mu \widehat{\Omega} - i\widehat{\Omega}^\dagger \partial_\mu \widehat{\Omega} \Big) \Big( g' B^\mu + g\widehat{\Omega}^\dagger A^\mu \widehat{\Omega} - i\widehat{\Omega}^\dagger \partial^\mu \widehat{\Omega} \Big) \Big\} \Big). \end{aligned} \tag{15}$$

But other terms are possible [5], as, for instance, from eq. (10)

$$S = \int d^4x \Big( -\frac{1}{4} G_{a\mu\nu} G_a^{\mu\nu} - \frac{1}{4} F_{\mu\nu} F^{\mu\nu}$$
$$+ \frac{m^2}{g^2} \left\{ v^\dagger \left( g' B_\mu + g \widehat{\Omega}^\dagger A_\mu \widehat{\Omega} - i \widehat{\Omega}^\dagger \partial_\mu \widehat{\Omega} \right) v \right\}^2 \Big). \tag{16}$$

Clearly the Stückelberg theory based on the gauge transformations in eq. (12) is not a good candidate as a limit of the ew model. First $\widehat{\Omega}$ contains a $U(1)$ factor which is really not required. Second, due to the freedom of introducing more terms bilinear in the vector fields, the Stückelberg mass formulation cannot reproduce, even at the tree level, the phenomenology electroweak interactions (as the neutral currents with the correct couplings or the $\rho$ parameter).

# 3 Extension to $SU(2) \otimes SU(2)$

Let us consider again the action (5). The form of the mass term suggests that one can enlarge the symmetry of the action by considering the right $SU(2)$ global transformations. This model has been studied at length by Bardeen and Shizuya in Ref. [3]. One can extend the symmetry to the right $SU(2)$ local transformations by introducing a new field such that

$$\begin{aligned} B'_\mu &= U_R B_\mu U_R^\dagger - \frac{i}{g'} \partial_\mu U_R U_R^\dagger \\ A'_\mu &= U_L A_\mu U_L^\dagger - \frac{i}{g} U_L \partial_\mu U_L^\dagger \\ \Omega' &= U_L \Omega U_R^\dagger. \end{aligned} \tag{17}$$

Then under local $SU_L(2) \otimes SU_R(2)$ transformations we have

$$\left( g' B_\mu + g \Omega^\dagger A^\mu \Omega - i \Omega^\dagger \partial^\mu \Omega \right) \to U_R \left( g' B_\mu + g \Omega^\dagger A^\mu \Omega - i \Omega^\dagger \partial^\mu \Omega \right) U_R^\dagger. \tag{18}$$

Consequently the only term invariant under $SU_L(2) \otimes SU_R(2)$ local transformations, that is bilinear in the gauge fields, is

$$\int d^4x \frac{m^2}{g^2}$$
$$Tr \ \left\{ \left( g' B_\mu + g \Omega^\dagger A_\mu \Omega - i \Omega^\dagger \partial_\mu \Omega \right) \left( g' B_\mu + g \Omega^\dagger A^\mu \Omega - i \Omega^\dagger \partial^\mu \Omega \right) \right\}. \tag{19}$$

We can add chiral fermions to this theory. The necessity to introduce a Yukawa coupling

$$\bar{\psi}_L \Omega \psi_R + \bar{\psi}_R \Omega^\dagger \psi_L \tag{20}$$

determines the transformation properties of the fermion fields and therefore the kinetic part of the action

$$\int d^4x \Big[ \bar{\psi}_L \gamma^\mu \left( i\partial_\mu + gA_\mu \right) \psi_L + \bar{\psi}_R \gamma^\mu \left( i\partial_\mu + g' B_\mu \right) \psi_R \Big]. \tag{21}$$

This model has three massive and three massless vector mesons at the tree level of the perturbation theory. In order to get a reasonable model, one has to devise a mechanism that removes two of the massless vector mesons. Moreover the correct quantum numbers of quark and leptons can be obtained by introducing a novel $U(1)_{L-B}$ symmetry in analogy to celebrated left-right symmetric models [6].

In this line of thought a Higgs field in the adjoint representation of $SU_R(2)$ can be introduced in order to generate a mass term for two of the three vector mesons $B^a_\mu$, keeping the third one massless. The massive $SU_R(2)$-gauge bosons as well as the adjoint Higgs field are to be made sufficiently heavy in order to reproduce the low-energy Standard Model phenomenology.

In a somehow different fashion an extended symmetry content can be obtained by gauging the hidden symmetry of non-linear $\sigma$ model [18]. This results in an extra local $SU(2)$ invariance, which is implemented by means of an additional set of heavy gauge bosons. Along these lines a phenomenologically viable extension of the SM has been derived in Refs. [18].

# 4 $SU(2) \otimes U(1)$ Symmetry

The toy model of the Section 3 suggests a way to build a Stückelberg theory of the ew model which is not in contrast with the gross features of phenomenology. Let us consider the reduction of the full symmetry $SU_L(2) \otimes SU_R(2)$ to $SU_L(2) \otimes U(1)$. When we consider this reduction on the Fermion sector one cannot forget that left-right symmetric models necessitate of an extra $U(1)$ invariance in order to distinguish leptons from quarks, as mentioned at the end of Section 3. Then the process is not simply the reduction of $SU_R(2) \to U_R(1)$. However when we consider the gauge sector by itself, then this reduction is achieved simply by $SU_R(2) \to U_R(1)$ i.e. by imposing invariance under the transformations (see eq. (17))

$$\begin{aligned} \tau_3 B'_\mu &= \tau_3 B_\mu - \frac{i}{g'} \partial_\mu U_R U_R^\dagger \\ A'_\mu &= U_L A_\mu U_L^\dagger - \frac{i}{g} U_L \partial_\mu U_L^\dagger \\ \Omega' &= U_L \Omega U_R^\dagger, \end{aligned} \tag{22}$$

with

$$U_L \in SU_L(2)$$

$$U_R = \exp(i\lambda\frac{\tau_3}{2}) \in U(1). \tag{23}$$

A reduction of the symmetry of the action allows to have more invariant terms. We use a fixed vector, eigenvector of $\tau_3$

$$v_+ = \begin{pmatrix} 1 \\ 0 \end{pmatrix} \tag{24}$$

then we have an invariant mass term described by the expression

$$\begin{aligned} & v_+^\dagger \ \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega\right) \left(g' B^\mu \frac{\tau_3}{2} + g\Omega^\dagger A^\mu \Omega - i\Omega^\dagger \partial^\mu \Omega\right) v_+ \\ & = \Phi^\dagger \left\{ \left(\frac{g'}{2} B_\mu + g A_\mu + i \overleftarrow{\partial}_\mu\right) \left(\frac{g'}{2} B^\mu + g A^\mu - i \overrightarrow{\partial^\mu}\right) \right\} \Phi \end{aligned} \tag{25}$$

where

$$\Phi = \Omega v_+, \qquad \Phi^\dagger \Phi = 1, \tag{26}$$

i.e. the formal limit of $\lambda \to \infty$ in the electroweak model.

In terms of scalar fields

$$\Phi = \begin{pmatrix} \phi_0 + i\phi_3 \\ i\phi_1 - \phi_2 \end{pmatrix} \tag{27}$$

An other invariant is built with

$$v_- = \begin{pmatrix} 0 \\ 1 \end{pmatrix} \tag{28}$$

and one gets

$$\begin{aligned} & v_-^\dagger \ \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega\right) \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A^\mu \Omega - i\Omega^\dagger \partial^\mu \Omega\right) v_- \\ & = \tilde{\Phi}^\dagger \left\{ \left(-\frac{g'}{2} B_\mu + g A_\mu + i \overleftarrow{\partial}_\mu\right) \left(-\frac{g'}{2} B^\mu + g A^\mu - i \overrightarrow{\partial^\mu}\right) \right\} \tilde{\Phi} \end{aligned} \tag{29}$$

where

$$\tilde{\Phi} \equiv \Omega v_- = \begin{pmatrix} i\phi_1 + \phi_2 \\ \phi_0 - i\phi_3 \end{pmatrix}. \tag{30}$$

Since

$$\tilde{\Phi} = \epsilon \Phi^* \tag{31}$$

where

$$\epsilon = \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}, \tag{32}$$

it is easy to prove that the two invariants in eqs. (25) and (29) are the same. Thus we can write eq. (25) also in the form

$$\frac{1}{2}\ Tr\ \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega\right) \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A^\mu \Omega - i\Omega^\dagger \partial^\mu \Omega\right). \tag{33}$$

This form will be useful to discuss custodial symmetry in section 5.

### 4.1 Mass terms

Now we consider the Yukawa couplings for the Fermi-Dirac fields. Since we have reduced the right symmetry $SU_R(2) \to U(1)$ we can build more invariants as in eq. (20). Then the Yukawa sector can be enlarged to a two-parameter space by

$$f\bar{u}_R\Phi^\dagger\psi_L + \tilde{f}\bar{d}_R\tilde{\Phi}^\dagger\psi_L + h.c. \tag{34}$$

which can be extended to the Cabibbo-Kobayashi-Maskawa [7] theory by considering the adequate number of Fermion families. The requirement that the Yukawa term in eq. (34) is invariant under $SU_L(2)\otimes U(1)$ leaves enough freedom to reproduce the standard internal quantum numbers of leptons and quarks. The only constraint on the $U(1)$ charge of the Fermions comes from the balance implied by the equation (34).

## 5 Custodial symmetry and Slavnov-Taylor identities

The fundamental problem of the existence of a theory which is the limit for infinite Higgs-self-coupling cannot avoid to exploit some general properties of the conventional ew model. Two items are exceedingly important for somewhat different reasons. We only sketch them here.

### 5.1 Custodial symmetry

The Standard model has a custodial symmetry [8] which prevents large corrections to the $\rho$ parameter [4]. The ew model becomes invariant under a *global* $SU_R(2)$ symmetry when the $U(1)$ coupling constant is switched off and the Yukawa couplings are put equal within a single flavor family. The Stückelberg formulation we are presenting justs hands on this symmetry.

This feature can be seen from eq. (33) in Section 4, where for $g' = 0$ and equal Yukawa couplings this symmetry is explicitly displayed. The consequences of this symmetry are similar to those of the ew model.

By changing basis of the global $SU_L(2)\otimes SU_R(2)$ transformations we introduce $SU_V(2)$ and $SU_A(2)$ transformations. In particular on the Stückelberg field we get

$$\begin{aligned} \delta_V\phi_0 &= 0 \\ \delta_V\phi_a &= \frac{\delta\omega_{Vc}}{2}\epsilon_{abc}\phi_b \end{aligned} \tag{35}$$

and

$$\begin{aligned} \delta_A\phi_0 &= -\frac{\delta\omega_{Aa}}{2}\phi_a \\ \delta_A\phi_a &= \frac{\delta\omega_{Aa}}{2}\phi_0. \end{aligned} \tag{36}$$

Eq.(35) corresponds to the choice $U_L = U_R$ in eq.(17), eq.(36) to $U_L = U_R^\dagger$. Setting $\Omega_{\rm Higgs} = \phi_0\ 1 + i\phi^a\tau^a$ after spontaneous symmetry breaking one has $< \Omega_{\rm Higgs} >= v\ 1$. $< \Omega_{\rm Higgs} >$ is left invariant under the transformation in eq.(17) provided that $U_L = U_R$. This choice gives rise to the custodial symmetry in eq.(35), leaving the field $\phi_0$ is invariant. The spontaneous breakdown does not affect the generators of the $SU_V(2)$ group of transformations.

### 5.2 Slavnov-Taylor identities

The BRST [9] invariance properties of the Stückelberg formulation of the electroweak model are identical to those of the standard model. The fact that $\phi_0$ is a composite field does not change its transformation properties, since $\Omega^\dagger\Omega = 1$ is a gauge- (and thus BRST-)invariant constraint. Thus we have for the $SU(2) \otimes U(1)$ sector

$$
\begin{array}{ll}
s_1\, A_{a\mu} = \partial_\mu c_a - A_{c\mu}\epsilon_{abc}c_b & s_0\, A_{a\mu} = 0 \\
s_1\, B_\mu = 0 & s_0\, B_\mu = \partial_\mu c_0 \\
s_1\, c_a = -\frac{1}{2}\epsilon_{abc}c_b c_c & s_0\, c_a = 0 \\
s_1\, c_0 = 0 & s_0\, c_0 = 0 \\
s_1\, \bar{c}_a = b_a & s_0\, \bar{c}_a = 0 \\
s_1\, b_a = 0 & s_0\, b_a = 0 \\
s_1\, \bar{c}_0 = 0 & s_0\, \bar{c}_0 = b_0 \\
s_1\, b_0 = 0 & s_0\, b_0 = 0 \\
s_1\, \phi_a = -\epsilon_{abc}c_b\phi_c + c_a\phi_0 & s_0\, \phi_1 = c_0\phi_2 \\
 & s_0\, \phi_2 = -c_0\phi_1 \\
 & s_0\, \phi_3 = -c_0\phi_0 \\
s_1\, \phi_0 = -c_a\phi_a & s_0\, \phi_0 = c_0\phi_3 \\
\text{.....matter part} &
\end{array}
\tag{37}
$$

where the fields $b_a,\ b_0$ are the Lagrange multipliers used to impose the Landau gauge. Clearly these BRST transformations lead to the same Slavnov Taylor identities [10] valid for the Standard Model in the Electroweak sector.

The composite nature of the field $\phi_0(x)$ is reflected in the formulation of the relevant Slavnov Taylor identities [10]. In order to define the correlation functions of $\phi_0(x)$ the latter has to be coupled in the tree-level approximation of the vertex functional $\Gamma^{(0)}$ to the external source $\beta(x)$. The external source $\beta^*(x)$ coupled to the BRST variation $s\phi_0(x)$ has also to be included. The $\beta, \beta^*$-dependence of $\Gamma^{(0)}$ is then

$$
\int d^4x \left(\beta(x)\phi_0(x) + \beta^*(x)s\phi_0(x)\right) . \tag{38}
$$

The Stückelberg model is not power-counting renormalizable. The cohomology of the relevant BRST differential in eq.(37) has been computed in [11]. It has been shown there that the most general deformation of the

action (in the space of local formal power series) is given, up to trivial invariants, by a strictly gauge invariant term plus winding number terms. The latter are irrelevant in perturbation theory. Moreover, there is no perturbative anomaly. As a consequence, the Stückelberg model turns out to be renormalizable in the modern sense of [12].

In order to discuss Physical Unitarity the construction of a nilpotent BRST charge $Q$ [13, 14] is required. The conditions on the operator $Q$ needed to establish Physical Unitarity have been given under fairly general assumptions (not restricted to the perturbative framework) in [15]. They provide constraints on the quantization procedures aiming at a non-perturbative definition of the Stückelberg model. Under the assumption that the subtraction scheme fulfills the ST identities, the cancellation mechanisms implied by perturbative Physical Unitarity have been analyzed in [16].

# 6 Other invariants

The Stückelberg mass term in eq. (33) has been constructed as an effective field theory in the spirit of mantaining most of the properties of the Standard model as it has been shown in Section 5. This requirement is important if one hopes to establish some relation between the Higgs formulation and the Stückelberg's one.

There is a further invariant under the *local* $SU(2) \otimes U(1)$ which is also a bilinear form in the vector fields [17]. This form can be constructed by noticing that

$$v_+^\dagger \ \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega\right) v_- \tag{39}$$

under the transformations in eq. (22) becomes

$$\exp(i\lambda)\ v_+^\dagger \ \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega\right) v_- \tag{40}$$

Then the following bilinear form in the gauge field is invariant

$$\begin{aligned}
& v_+^\dagger \ \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A_\mu \Omega - i\Omega^\dagger \partial_\mu \Omega\right) v_- \\
& v_-^\dagger \ \left(g' B_\mu \frac{\tau_3}{2} + g\Omega^\dagger A^\mu \Omega - i\Omega^\dagger \partial^\mu \Omega\right) v_+ \\
& = \left\{\Phi^\dagger \left(gA_\mu + i \overleftarrow{\partial}_\mu\right) \tilde{\Phi}\right\} \ \left\{\tilde{\Phi}^\dagger \left(gA^\mu - i \overrightarrow{\partial^\mu}\right) \Phi\right\}.
\end{aligned} \tag{41}$$

This term can be dismissed only on the basis of the requirement that a custodial symmetry is present as discussed in Section 5. In the standard model it is not present since it contains couplings that cause non renormalizability of the theory.

## 7 Conclusions

In the present paper we have shown that the formal limit of infinite Higgs potential ($\lambda \to \infty$) can be casted in a theory with a Stückelberg mass. Moreover, always at the formal level, the proposed limit enjoys the same custodial symmetry and the BRST invariance properties as the Standard Electroweak model. We make clear that our work does not prove the existence of such a limit and, if such a limit exists, that no Higgs boson is present. In fact a physical boson particle might show up in many ways, e.g. as a non perturbative effects. Nevertheless it is rather surprising that such a powerful tool as the Slavnov-Taylor identities can be traced also in the limit of infinite Higgs potential.

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